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Large CP Phases and the Cancellation Mechanism in EDMs in SUSY, String and Brane Models

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Abstract

We show that EDMs obey a simple approximate scaling under the transformation $m_0 \rightarrow \lambda m_0, m_{\frac{1}{2}} \rightarrow \lambda m_{\frac{1}{2}}$ in the large μ region when μ itself obeys the same scaling, ie., $\mu \rightarrow \lambda \mu$. In the scaling region the knowledge of a single point in the MSSM parameter space where the cancellation in the EDMs occur allows one to generate a trajectory in the $m_0 - m_{\frac{1}{2}}$ plane where the cancellation mechanism holds and the EDMs are small. We illustrate these results for MSSM with radiative electro-weak symmetry breaking constraints. We also discuss a class of D brane models based on Type IIB superstring compactifications which have non-universal phases in the gaugino mass sector and allow large CP violating phases consistent with the EDM constraints through the cancellation mechanism. The scaling in these D brane models and in a heterotic string model is also discussed.

1 Introduction

Supersymmetric theories contain new sources of CP violation which arise from the phases of the soft SUSY breaking parameters which are in general complex. The CP violating phases associated with the complex soft SUSY breaking parameters are typically large, i.e. $O(1)$, and pose a problem regarding the satisfaction of the current experimental limits on the neutron and on the electron EDM. For the neutron the current experimental limit is[1]

$$|d_n| < 6.3 \times 10^{-26} ecm \quad (1)$$

and for the electron the limit is[2]

$$|d_e| < 4.3 \times 10^{-27} ecm. \quad (2)$$

Various remedies have been suggested in the literature to overcome this problem. The first of these is the suggestion that the phases are small[3, 4] $O(10^{-2})$. However, small phases constitute a fine tuning and are thus undesirable. Another suggestion is that the sparticle mass spectrum is heavy in the several TeV range to suppress the EDMs[5]. A third possibility suggested is that there are internal cancellations among the various contributions to the neutron and to the electron EDM leading to compatibility with experiment with large phases and a SUSY spectrum that is still within the reach of the accelerators[6]. There have been further developments[7, 8, 9, 10, 11, 12] and applications of this idea to explore the effects of large CP violating phases on dark matter analyses[13, 14, 15, 16, 17], on $g_\mu - 2$ [18], and on other low energy phenomena[19, 20, 21, 22, 23, 24]. The focus of this paper is to show that in theories where the higgs mixing parameter μ obeys the simple scaling behavior as the rest of the SUSY masses the EDMs exhibit a simple scaling behavior under the simultaneous scaling on m_0 and $m_{\frac{1}{2}}$. The scaling property of EDMs allows one to promote a single point in the SUSY parameter space where cancellations occur to a trajectory in the $m_0 - m_{\frac{1}{2}}$ plane. The scaling phenomena also has implications for the satisfaction of the EDM constraints in string and D brane models. The outline of the paper is as follows: In Sec.2 we discuss the scaling transformations and the properties of the relevant SUSY spectrum under scaling in the region of large μ . In Sec.3 we discuss the properties of the EDMs under scaling in this region. In Sec.4 we discuss the algorithm for the satisfaction of the EDM constraints. We also investigate the parameter space where μ is large and show that in this region scaling can be used to generate

trajectories in the $m_0 - m_{\frac{1}{2}}$ plane where the cancellation mechanism holds. The cancellation mechanism in string models and D brane models is discussed in Sec.5. Conclusions are given in Sec.6.

2 Scaling

In this section we discuss the properties of the chargino and the neutralino mass eigen-values and eigen vectors under the scale transformation

$$m_0 \rightarrow \lambda m_0, \quad m_{\frac{1}{2}} \rightarrow \lambda m_{\frac{1}{2}} \quad (3)$$

In general the eigen-spectrum will have no simple property under this transformation since the chargino and the neutralino mass matrices contain non-scaling parameters M_W and M_Z . However, simple scaling properties emerge when $|\mu| \gg M_Z$. In MSSM μ is an independent parameter and has no scaling property under Eq.(3). However, in scenarios with radiative breaking of the electro-weak symmetry μ is determined via one of the extrema equations by varying the effective potential

$$\mu^2 = \frac{1}{2}M_Z^2 + \frac{\tilde{m}_{H1}^2 - \tilde{m}_{H1}^2 \tan^2 \beta}{\tan^2 \beta - 1} \quad (4)$$

where $\tilde{m}_{H_i}^2 = m_{H_i}^2 + \Sigma_i$ (i=1,2) and where Σ_i is the one loop correction to the Higgs mass. In the limit $|\mu| \gg M_Z$, μ^2 becomes a homogeneous polynomial of degree 2 in m_0 and $m_{\frac{1}{2}}$, and thus under the transformation of Eq.(3) it has the property

$$\mu \rightarrow \lambda \mu \quad (5)$$

From now on we shall consider the class of models where Eq.(5) holds. Next let us consider the chargino mass matrix with the most general set of phases

$$M_C = \begin{pmatrix} |\tilde{m}_2| e^{i\xi_2} & \sqrt{2}m_W \sin \beta e^{-i\chi_2} \\ \sqrt{2}m_W \cos \beta e^{-i\chi_1} & |\mu| e^{i\theta_\mu} \end{pmatrix} \quad (6)$$

where our notation is as in Ref.[6]. By the transformation $M_C = B_R M'_C B_L^\dagger$, where $B_R = \text{diag}(e^{i\xi_2}, e^{-i\chi_1})$ and $B_L = \text{diag}(1, e^{i(\chi_2 + \xi_2)})$, the chargino mass matrix can be written in the form

$$M'_C = \begin{pmatrix} |\tilde{m}_2| & \sqrt{2}m_W \sin \beta \\ \sqrt{2}m_W \cos \beta & |\mu| e^{i(\tilde{\theta})} \end{pmatrix} \quad (7)$$

where $\tilde{\theta} = \theta_\mu + \xi_2 + \chi_1 + \chi_2$. We can diagonalize the matrix M'_C by the biunitary transformation

$$U_R'^\dagger M'_C U_L = \text{diag}(|\tilde{m}_{\chi_1^+}| e^{i\gamma_1}, |\tilde{m}_{\chi_2^+}| e^{i\gamma_2}) \quad (8)$$

In the limit of $|\mu| > (M_W, |\tilde{m}_2|)$ the eigen-values of the chargino mass matrix are[25]

$$\tilde{m}_{\chi_1^+} \simeq |\mu|, \quad \tilde{m}_{\chi_2^+} \simeq |\tilde{m}_2| \quad (9)$$

These relations were derived originally in the absence of CP violating phases in the limit of large μ in supergravity models with radiatively induced breaking of the electro-weak symmetry. Here we find that the relations continue to hold when CP violating phases are included. The matrices U'_R and U_L in the large μ limit may be expanded as follows

$$\begin{aligned} U_L &= 1 + U_L^{(1)} \left(\frac{M_W}{|\mu|} \right) + U_L^{(2)} \left(\frac{M_W^2}{|\mu|^2} \right) + \dots \\ U'_R &= 1 + U_R^{(1)} \left(\frac{M_W}{|\mu|} \right) + U_R^{(2)} \left(\frac{M_W^2}{|\mu|^2} \right) + \dots \end{aligned} \quad (10)$$

where $U_{L,R}^{(1,2)}$ are scale independent matrices and are given by

$$\begin{aligned} U_L^{(1)} &= \begin{pmatrix} 0 & \sqrt{2} \cos \beta e^{i\tilde{\theta}} \\ -\sqrt{2} \cos \beta e^{-i\tilde{\theta}} & 0 \end{pmatrix} \\ U_L^{(2)} &= \begin{pmatrix} -\cos^2 \beta & 0 \\ 0 & -\cos^2 \beta \end{pmatrix} \\ U_R^{(1)} &= \begin{pmatrix} 0 & \sqrt{2} \sin \beta e^{-i\tilde{\theta}} \\ -\sqrt{2} \sin \beta e^{i\tilde{\theta}} & 0 \end{pmatrix} \\ U_R^{(2)} &= \begin{pmatrix} -\sin^2 \beta & 0 \\ 0 & -\sin^2 \beta \end{pmatrix} \end{aligned} \quad (11)$$

By defining $U_R = U'_R \times \text{diag}(e^{-i\gamma_1}, e^{-i\gamma_2})$ one can have

$$U_R^\dagger M'_C U_L = \text{diag}(|\tilde{m}_{\chi_1^+}|, |\tilde{m}_{\chi_2^+}|) \quad (12)$$

Thus to the leading order under the transformation of Eqs.(3) and (5) one has

$$|m_{\chi_i^+}| \rightarrow \lambda |m_{\chi_i^+}|, \quad i = 1, 2 \quad (13)$$

and the relevant matrix elements of the EDMs will have the following scale transformations:

$$\begin{aligned} \text{Im}(U_{L2i} U_{R1i}^*) &\rightarrow \frac{1}{\lambda} \text{Im}(U_{L2i} U_{R1i}^*) \\ \text{Im}(U_{L1i} U_{R2i}^*) &\rightarrow \frac{1}{\lambda} \text{Im}(U_{L1i} U_{R2i}^*) \end{aligned} \quad (14)$$

We discuss now the neutralino mass matrix

$$\begin{pmatrix} |\tilde{m}_1|e^{i\xi_1} & 0 & -M_z \sin \theta_W \cos \beta e^{-i\chi_1} & M_z \sin \theta_W \sin \beta e^{-i\chi_2} \\ 0 & |\tilde{m}_2|e^{i\xi_2} & M_z \cos \theta_W \cos \beta e^{-i\chi_1} & -M_z \cos \theta_W \sin \beta e^{-i\chi_2} \\ -M_z \sin \theta_W \cos \beta e^{-i\chi_1} & M_z \cos \theta_W \cos \beta e^{-i\chi_2} & 0 & -|\mu|e^{i\theta_\mu} \\ M_z \sin \theta_W \sin \beta e^{-i\chi_1} & -M_z \cos \theta_W \sin \beta e^{-i\chi_2} & -|\mu|e^{i\theta_\mu} & 0 \end{pmatrix}. \quad (15)$$

We define the matrix X that diagonalizes M_{χ^0} so that

$$X^T M_{\chi^0} X = \text{diag}(\tilde{m}_{\chi_1^0}, \tilde{m}_{\chi_2^0}, \tilde{m}_{\chi_3^0}, \tilde{m}_{\chi_4^0}) \quad (16)$$

In the limit $|\mu| > \{M_Z, |\tilde{m}_1|, |\tilde{m}_2|\}$ the neutralino mass eigen-values have the following form

$$\tilde{m}_{\chi_1^0} \simeq |\tilde{m}_1|, \tilde{m}_{\chi_2^0} \simeq |\tilde{m}_2|, \tilde{m}_{\chi_3^0} \simeq |\mu|, \tilde{m}_{\chi_4^0} \simeq |\mu| \quad (17)$$

Again the scaling relations of Eq.(17) were originally derived in the limit of large μ and no CP phases and our analysis shows that these relations continue to hold when large CP violating phases are included. From Eq.(17) we find that in the large μ limit under the transformations of Eqs.(3) and (5) one has

$$m_{\chi_i^0} \rightarrow \lambda m_{\chi_i^0} \quad (i = 1 - 4) \quad (18)$$

In the large μ limit the diagonalizing matrix X has the expansion

$$X = X^{(0)} + X^{(1)}\left(\frac{M_Z}{|\mu|}\right) + O\left(\frac{M_Z^2}{|\mu|^2}\right) \quad (19)$$

where $X^{(0),(1)}$ are scale independent matrices. Now we discuss the behavior of the diagonalizing matrix D of the sfermion $(mass)^2$ matrix under the scaling transformations where

$$D^\dagger M_{\tilde{f}}^2 D = \text{diag}(M_{\tilde{f}_1}^2, M_{\tilde{f}_2}^2) \quad (20)$$

For light flavors the scale transformations for the mass eigen states are

$$M_{\tilde{f}_i} \rightarrow \lambda M_{\tilde{f}_i} \quad (i = 1 - 2) \quad (21)$$

and the matrix elements of D have the following transformations under the scaling transformation of Eqs.(3) and (5): $D_{11}, D_{22} \rightarrow D_{11}, D_{22}$; $D_{12}, D_{21} \rightarrow \frac{1}{\lambda} D_{12}, D_{21}$. We note, however, that for the light flavors (the electron, the up quark and the down quark) one has $|D_{12}, D_{21}| < |D_{11}, D_{22}|$. For the heavy flavors (i.e., the top and the bottom quarks) which are relevant to the six dimensional purely gluonic operator, the behavior of the eigen values and of the diagonalizing matrices are much more complicated and will be discussed later.

3 Scaling Properties of EDMs

In the analysis below we shall use the notation of Refs.[6]. However, we will make the notation explicit where necessary. The chargino contribution to the EDM of the up quark is given by

$$d_{u\text{-chargino}}^E/e = \frac{-\alpha_{EM}}{4\pi \sin^2 \theta_W} \sum_{k=1}^2 \sum_{i=1}^2 \text{Im}(\Gamma_{uik}) \frac{\tilde{m}_{\chi_i^+}}{M_{dk}^2} [Q_{\tilde{d}} \text{B}(\frac{\tilde{m}_{\chi_i^+}^2}{M_{dk}^2}) + (Q_u - Q_{\tilde{d}}) \text{A}(\frac{\tilde{m}_{\chi_i^+}^2}{M_{dk}^2})], \quad (22)$$

where $A(r) = (2(1-r))^{-2}(3-r+2\ln r(1-r)^{-1})$ and

$$\Gamma_{uik} = \kappa_u V_{i2}^* D_{d1k} (U_{i1}^* D_{d1k}^* - \kappa_d U_{i2}^* D_{d2k}^*) \quad (23)$$

and $\kappa_u = m_u e^{-i\chi_2}/\sqrt{2}M_W \sin \beta$. Because of the smallness of m_u , we can ignore the second part of Γ_{uik} and the bigger component of it would be that of $k=1$ and it could be written in terms of $U_{L,R}$ as

$$\Gamma_{ui1} \simeq |\kappa_u| |D_{d11}|^2 U_{L2i} U_{R1i}^* \quad (24)$$

which under the scaling transformation behaves as

$$\Gamma_{ui1} \rightarrow \frac{1}{\lambda} \Gamma_{ui1} \quad (25)$$

So the chargino component of the electric operator for the up quark $d_u^{\chi+}$ has the scale transformation

$$d_u^{\chi+} \rightarrow \frac{1}{\lambda^2} d_u^{\chi+} \quad (26)$$

and the same transformation holds for the down quark and for the electron

$$d_{d,e}^{\chi+} \rightarrow \frac{1}{\lambda^2} d_{d,e}^{\chi+} \quad (27)$$

The neutralino exchange contribution to a fermion is given by[6]

$$d_{f\text{-neutralino}}^E/e = \frac{\alpha_{EM}}{4\pi \sin^2 \theta_W} \sum_{k=1}^2 \sum_{i=1}^4 \text{Im}(\eta_{fik}) \frac{\tilde{m}_{\chi_i^0}}{M_{fk}^2} Q_{\tilde{f}} \text{B}(\frac{\tilde{m}_{\chi_i^0}^2}{M_{fk}^2}) \quad (28)$$

where

$$\eta_{fik} = (a_0 X_{1i} D_{f1k}^* + b_0 X_{2i} D_{f1k}^* - \kappa_f X_{bi} D_{f2k}^*) (c_0 X_{1i} D_{f2k} - \kappa_f X_{bi} D_{f1k}) \quad (29)$$

Here $b=3(4)$ for $T_{3q} = -\frac{1}{2}(\frac{1}{2})$, $a_0 = -\sqrt{2} \tan \theta_W (Q_f - T_{3f})$, $b_0 = -\sqrt{2} T_{3f}$, $c_0 = \sqrt{2} \tan \theta_W Q_f$. κ_u is defined following Eq.(23) and $\kappa_{d,e}$ is given by $\kappa_{d,e} = m_{d,e} e^{-i\chi_1}/\sqrt{2}M_W \cos \beta$. Because of the smallness of κ_f one can write η_{fik} as

$$\eta_{fik} \simeq a_0 c_0 X_{1i}^2 D_{f1k}^* D_{f2k} + b_0 c_0 X_{1i} X_{2i} D_{f1k}^* D_{f2k} \quad (30)$$

and by using the expansion of the matrix X of Eq.(19) one can write

$$\eta_{fik} \simeq a_0 c_0 X_{1i}^{(0)2} D_{f1k}^* D_{f2k} + b_0 c_0 X_{1i}^{(0)} X_{2i}^{(0)} D_{f1k}^* D_{f2k} \quad (31)$$

Since the transformation for $D_{f1k}^* D_{f2k}$ for $k = 1, 2$ is given by

$$D_{f1k}^* D_{f2k} \rightarrow \frac{1}{\lambda} D_{f1k}^* D_{f2k} \quad (32)$$

the neutralino contribution for the electric operator for both the quarks and the leptons behaves as:

$$d_{f-\text{neutralino}}^E \rightarrow \frac{1}{\lambda^2} d_{f-\text{neutralino}}^E \quad (33)$$

Eqs.(27) and (33) imply that d_e satisfies the scaling property

$$d_e \rightarrow \frac{1}{\lambda^2} d_e \quad (34)$$

Next we discuss the EDM components for the quarks which contains the contributions from several operators, i.e., the electric dipole operator, the color dipole operator and the purely gluonic dimension six operator.

$$d_q = d_q^E + d_q^C + d_q^G \quad (35)$$

For the electric dipole the chargino and the neutralino contributions have already been discussed. For the gluino exchange contribution one has

$$d_{q-\text{gluino}}^E/e = \frac{-2\alpha_s}{3\pi} m_{\tilde{g}} Q_{\tilde{q}} \text{Im}(\Gamma_q^{11}) \left[\frac{1}{M_{\tilde{q}1}^2} \text{B}\left(\frac{m_{\tilde{g}}^2}{M_{\tilde{q}1}^2}\right) - \frac{1}{M_{\tilde{q}2}^2} \text{B}\left(\frac{m_{\tilde{g}}^2}{M_{\tilde{q}2}^2}\right) \right]. \quad (36)$$

where $\Gamma_q^{1k} = e^{-i\xi_3} D_{q2k} D_{q1k}^*$, $\Gamma_q^{12} = -\Gamma_q^{11}$ and

$$\text{Im}(\Gamma_q^{11}) = \frac{m_q}{M_{\tilde{q}1}^2 - M_{\tilde{q}2}^2} (m_0 |A_q| \sin(\alpha_q - \xi_3) + |\mu| \sin(\theta_\mu + \chi_1 + \chi_2 + \xi_3) |R_q|), \quad (37)$$

In the $|\mu|/M_Z \gg 1$ limit we find that $\text{Im}(\Gamma_q^{11})$ scales as $1/\lambda$ under the scaling of Eq.(3) and $d_{q-\text{gluino}}^E$ exhibits the same scaling behavior, i.e., $d_{q-\text{gluino}}^E \rightarrow \frac{1}{\lambda^2} d_{q-\text{gluino}}^E$.

Next we consider the chromoelectric dipole moment \tilde{d}^C contribution to the quark EDM. It is given by

$$d_q^C = \frac{e}{4\pi} \tilde{d}_q^C \eta^c \quad (38)$$

where η^c is the renormalization group evolution of the chromo-electric operator from the electro-weak scale to the hadronic scale and numerically $\eta^c \sim 3.3$. Contributions to \tilde{d}_q^C arise from the gluino, from the chargino and from the neutralino exchanges and we reproduce here the analytic expressions derived in Ref.[6].

$$\tilde{d}_{q-\text{gluino}}^C = \frac{g_s \alpha_s}{4\pi} \sum_{k=1}^2 \text{Im}(\Gamma_q^{1k}) \frac{m_{\tilde{g}}}{M_{\tilde{q}k}^2} \text{C}\left(\frac{m_{\tilde{g}}^2}{M_{\tilde{q}k}^2}\right), \quad (39)$$

$$\tilde{d}_{q\text{-chargino}}^C = \frac{-g^2 g_s}{16\pi^2} \sum_{k=1}^2 \sum_{i=1}^2 \text{Im}(\Gamma_{qik}) \frac{\tilde{m}_{\chi_i^+}}{M_{\tilde{q}k}^2} \text{B}\left(\frac{\tilde{m}_{\chi_i^+}^2}{M_{\tilde{q}k}^2}\right), \quad (40)$$

and

$$\tilde{d}_{q\text{-neutralino}}^C = \frac{g_s g^2}{16\pi^2} \sum_{k=1}^2 \sum_{i=1}^4 \text{Im}(\eta_{qik}) \frac{\tilde{m}_{\chi_i^0}}{M_{\tilde{q}k}^2} \text{B}\left(\frac{\tilde{m}_{\chi_i^0}^2}{M_{\tilde{q}k}^2}\right), \quad (41)$$

where our notation is as in Ref.[6]. The expansion of these contributions in the limit $|\mu|/M_Z \gg 1$ following the same procedure as for the electric dipole case shows that in this limit \tilde{d}^C again shows the scaling behavior $\tilde{d}^C \rightarrow \frac{1}{\lambda^2} \tilde{d}^C$ under the transformations of Eqs.(3) and (5). Finally, we consider the contribution of the purely gluonic dimension six operator. It is given by

$$d_q^G = \frac{eM}{4\pi} \tilde{d}_q^G \eta^G \quad (42)$$

where η^G is the renormalization group evolution of the purely gluonic dimension six operator from the electro-weak scale down to the hadronic scale ($\eta^G \simeq 3.3$) and M is the chiral symmetry breaking scale ($M \simeq 1.19$ GeV) and \tilde{d}_q^G is given by[4]

$$\tilde{d}_q^G = -3\alpha_s \left(\frac{g_s}{4\pi m_{\tilde{g}}}\right)^3 (m_t(z_1^t - z_2^t) \text{Im}(\Gamma_t^{12}) H(z_1^t, z_2^t, z_t) + m_b(z_1^b - z_2^b) \text{Im}(\Gamma_b^{12}) H(z_1^b, z_2^b, z_b)) \quad (43)$$

where

$$\Gamma_q^{1k} = e^{-i\xi_3} D_{q2k} D_{q1k}^*, z_\alpha^q = \left(\frac{M_{\tilde{q}\alpha}}{m_{\tilde{g}}}\right)^2, z_q = \left(\frac{m_q}{m_{\tilde{g}}}\right)^2 \quad (44)$$

The behavior of Γ_t^{12} , z_α^q and z_q under the scaling transformation is a complicated one because of the largeness of the quark masses involved and even if we were in a region where one can ignore these masses compared to the other mass scales in the problem one finds that the behavior of d^G is different from that of the other components i.e. $d^G \rightarrow \frac{1}{\lambda^4} d^G$. Thus the scaling property of d_q will be more complicated. However, as λ gets large the contribution of d_q^G will fall off faster than the contribution of d_q^E and d_q^C and in this case one will have the scaling $d_q \rightarrow \frac{1}{\lambda^2} d_q$ and so also the neutron edm d_n will behave as

$$d_n \rightarrow \frac{1}{\lambda^2} d_n \quad (45)$$

We note, however, that the question of how soon the scaling sets in as we scale in λ depends on the part of the parameter space one is in.

4 Satisfaction of EDM Constraints

In the work of Ref.[6] it was shown that the quark and the lepton EDMs in general depend on ten independent phases which were classified there providing one with considerable freedom for the satisfaction of the EDM constraints. Numerical analyses show the existence of significant regions of the parameter space where the cancellation mechanism holds. We describe below a straightforward technique for accomplishing the satisfaction of the EDM constraints. These techniques are already well understood and we codify them here for the benefit of the reader. For the case of the electron one finds that the chargino component of the electron is independent of ξ_1 and the electron EDM as a whole is independent of ξ_3 . Thus the algorithm to discover a point of simultaneous cancellation for the electron EDM and for the neutron EDM is a straightforward one. For a given set of parameters except ξ_1 we start varying ξ_1 till we reach the cancellation for the electron EDM since only one of its components (the neutralino) is affected by that parameter. Once the electric dipole moment constraint on the electron is satisfied we vary ξ_3 which affects only the neutron edm keeping all other parameters fixed. By using this simple algorithm one can generate any number of simultaneous cancellations. In the numerical analysis of the EDMs we also take into account the two loop diagrams of the type discussed in Ref.[26]. However, we find that in the small $\tan\beta$ region these diagrams do not make any substantial contributions to the EDMs.

We discuss now the lepton and the neutron EDMs in the region where the scaling relation on the lepton and the neutron EDMs of Eqs.(34) and (45) hold. Suppose we have a point in the parameter space where the lepton and the quark EDMs vanish, i.e., $d_e = 0, d_q = 0$. The interesting observation is that this cancellation constraint is preserved under scaling provided one is in the scaling region, i.e., Eqs.(34) and (45) hold. Thus given a point in the parameter space where cancellations occur one can generate a trajectory in the $m_0 - m_{\frac{1}{2}}$ plane by a simple scaling of m_0 and $m_{\frac{1}{2}}$ using Eqs.(3) and (5). In practice the cancellation is not designed to be perfect and the scaling properties of d_e given by Eq.(34) and of d_n given by Eq.(45) are only approximate. Thus under the scaling transformation some minor adjustment of the other parameters will in general be necessary. The length of the trajectory depends on the part of the parameter space one is in. For some cases it is found that the trajectory can be long enough to cover the range of the parameter space consistent with naturalness. An example of this phenomenon is shown in Fig.1 where five trajectories are generated, and where each trajectory

is generated from a single cancellation point for low values of m_0 and $m_{\frac{1}{2}}$ by simple scaling. We notice, however, that there is an empty region in trajectory 5 where the cancellation under scaling does not hold. However, we have checked that with a very minor adjustments in the values of the other parameters we can restore the cancellation. Thus each of the trajectories satisfy the EDM constraints with the values of A_0 , $\tan\beta$, and phase angles fixed as we move along the trajectory. As we move on the trajectory to the higher mass regions we have a natural suppression besides the cancellation suppression. However, the cancellation is still necessary except for the extreme ends of each trajectory. In Fig.2 we exhibit the EDM of the neutron corresponding to the five trajectories of Fig.1. We find that all the trajectories are consistent with the current experimental constraint on the neutron EDM. In Fig.3 we plot the EDM of the electron corresponding to the five trajectories of Fig.1. Again we find that all the trajectories are consistent with the current experimental constraint on the electron EDM.

In summary a convenient procedure for generating a trajectory in the $m_0 - m_{\frac{1}{2}}$ plane where cancellations of the EDMs occur, consists of finding a single point in the MSSM parameter space with low values of m_0 and $m_{\frac{1}{2}}$ under the constraint of the radiative breaking of the electro-weak symmetry using the algorithm described in the beginning of this section where the cancellation in EDMs of the electron and of the neutron occur consistent with Eqs.(1) and (2). One then computes the EDMs using Eqs.(3) and (5) for $\lambda > 1$ and typically one finds that the EDM constraints are maintained with only minor adjustment of other parameters. The onset of the scaling behavior itself will depend on the values of the other MSSM parameters. We emphasize that in some cases the subleading terms in the scaling law may be significant and could generate new cancellations if they change sign as we scale upward in λ . While such points violate the scaling law, they are nonetheless acceptable since there is an even greater satisfaction of the EDM constraints for this case. In Figs. 4 and 5 we plot $\log_{10}\lambda^2|d_{e,n}|$ as a function of $m_{1/2}$ and we see support of the scaling idea here. It is important to keep in mind that the method we outlined here is only an approximation and should be used keeping that in mind. The method would work best if one is in the scaling region or close to it. Certainly it should be of relevance in exploring at least a part of the parameter space where these conditions are met.

5 String and Brane Models and EDM cancellations

We discuss now CP violation and cancellations in EDMs for the case of string and brane models. Recently, the progress in string dualities has led to the formulation of a new class of models based on M theory compactified on $CY \times S^1/Z_2$ and models in the framework of Type IIB orientifolds. We shall focus here on Type IIB orientifold models which have received significant attention recently[27]. Specifically we shall consider models with compactification of the Type IIB theory on a six-torus $T^6 = T^2 \times T^2 \times T^2$ of the type discussed in Ref.[28]. In scenarios of this type as in SUSY models and other string models additional sources of CP violation can arise through the breaking of supersymmetry. The mechanism of breaking of supersymmetry here is not fully understood. However, one can still make some progress by phenomenologically parametrizing how supersymmetry breaks. An efficient way of doing so is in terms of the VEVs of the dilaton field (S) and of the moduli fields T_i and for the case when the vacuum energy is set to zero one has that F type supersymmetry breaking may be parametrized by[28]

$$\begin{aligned} F^S &= \sqrt{3}m_{\frac{3}{2}}(S + S^*) \sin \theta e^{-i\gamma_S} \\ F^i &= \sqrt{3}m_{\frac{3}{2}}(T + T^*) \cos \theta \Theta_i e^{-i\gamma_i} \end{aligned} \quad (46)$$

where θ , Θ_i parametrize the Goldstino direction in the S , T_i field space and γ_S and γ_i are the F^S and F^i phases, and $\Theta_1^2 + \Theta_2^2 + \Theta_3^2 = 1$. The Type IIB compactified models of the type mentioned above contain 9 branes, 7_i ($i=1,2,3$) branes, 5_i ($i=1,2,3$) branes and 3 branes. N=1 supersymmetry constraints require that not all the branes can simultaneously be present, and thus one can have either 9 branes and 5_i branes or 7_i branes and 3 branes. Recently the work of Ref.[11] investigated the EDM constraints on models based on 5_i ($i=1,2$) branes which belong to the general class of models discussed in Ref.[28]. It was shown and that this model exhibits non-universalities in the phases of the gaugino masses and that cancellations in the EDMs arise and one can achieve satisfaction of the EDM constraints consistent with experiment[11, 12]. Our own analysis of this model further confirms the existence of the cancellations for the EDMs in the parameter space of this model.

We discuss here the models based on 9 branes and one from the set of 5_i branes which we choose to be 5_1 where the Standard Model gauge group is distributed between the two branes. Like the models based on 5_i ($i=1,2$) branes, these models

also contain non-universalities of the gaugino phases due to different gauge kinetic energy functions associated with 9 branes and 5_i branes, i.e., $f_9 = S$, and $f_{5_i} = T_i$. However, the nature of the soft SUSY breaking is different in these models from the ones based on 5_i branes. Thus it is interesting to investigate the question of large CP violating phases and of cancellations in the EDMs in this type of models. In models where more than one type of branes are involved the unification of the gauge couplings requires fine tuning. For the case of models based on the 9 brane and the 5_1 brane the unification of gauge couplings is more difficult than in the case when the gauge groups are embedded on two different same type branes. A full discussion of this topic is outside the scope of this work. However, we wish to note that contributions from extra matter and twisted moduli[28] could be important in a realistic analysis of the gauge coupling unification in this case. For the purpose of the analysis we shall simply assume that the unification does occur at the usual scale of $\sim 10^{16}$ GeV. We emphasize that the issue of cancellations in the EDMs is largely independent of the issue of the gauge coupling unification and thus the conclusions of our analysis are largely independent of this issue.

Below we consider the following two ways to embed the Standard Model gauge group on the 9 branes and 5_1 branes.

Case I:

Here we consider the possibility that the $SU(3)_C \times U(1)_Y$ is associated with the 9 brane and the $SU(2)_L$ is associated with the 5_1 brane. Further we assume that the $SU(2)_R$ singlet states are associated with the nine-brane sector, while the $SU(2)_L$ doublet states arise from the intersection of 9-brane and 5_1 -brane sector as in Case I. In this model we find using the general formulae of Ref.[28] the following results: the $SU(2)_R$ singlets have the common mass m_9 and the $SU(2)_L$ doublets have the common mass m_{95_1} where

$$m_9^2 = m_{\frac{3}{2}}^2 (1 - 3 \cos^2 \theta \Theta_1^2) \quad (47)$$

$$m_{95_1}^2 = m_{\frac{3}{2}}^2 (1 - \frac{3}{2} \cos^2 \theta (1 - \Theta_1^2)). \quad (48)$$

The $SU(3)$, $SU(2)$ and $U(1)$ gaugino masses \tilde{m}_i ($i=1,2,3$) are given by

$$\begin{aligned} \tilde{m}_1 &= \sqrt{3} m_{\frac{3}{2}} \sin \theta e^{-i\gamma_S} = \tilde{m}_3 = -A_0, \\ \tilde{m}_2 &= \sqrt{3} m_{\frac{3}{2}} \cos \theta \Theta_1 e^{-i\gamma_1} \end{aligned} \quad (49)$$

In the analysis of the EDMs we shall treat the phase of μ to be a free parameter and the magnitude of μ is determined by the radiative breaking of the electro-weak

symmetry. In order to avoid tachyons we impose the constraint $\cos^2\theta\Theta_1^2 < 1/3$. In Fig.6 we exhibit the cancellation phenomenon for the EDMs for this case in the presence of large CP violating phases.

Case II:

The second possibility is that the $SU(3)_C \times U(1)_Y$ is associated with the 5_1 brane and the $SU(2)_L$ is associated with the 9 brane. Regarding the matter fields we assume that the $SU(2)_R$ singlet states are associated with the 5_1 sector, while the $SU(2)_L$ doublet states arise from the intersection of the 5_1 -brane and the 9-brane sector. Although this case is T dual to Case I the pattern of soft masses is different after the breaking of supersymmetry in the two cases. Thus after SUSY breaking one finds here that the $SU(2)_R$ singlet masses have the common mass m_{5_1} and the $SU(2)_L$ doublet masses have the common mass m_{95_1} where

$$m_{5_1}^2 = m_{\frac{3}{2}}^2(1 - 3\sin^2\theta) \quad (50)$$

$$m_{95_1}^2 = m_{\frac{3}{2}}^2(1 - \frac{3}{2}\cos^2\theta(1 - \Theta_1^2)) \quad (51)$$

while the $SU(3)$, $SU(2)$ and $U(1)$ gaugino masses are given by

$$\begin{aligned} \tilde{m}_1 &= \sqrt{3}m_{\frac{3}{2}}\cos\theta\Theta_1e^{-i\gamma_1} = \tilde{m}_3 = -A_0, \\ \tilde{m}_2 &= \sqrt{3}m_{\frac{3}{2}}\sin\theta e^{-i\gamma_S} \end{aligned} \quad (52)$$

To guarantee that there are no tachyons we impose the constraint $\sin^2\theta < 1/3$. We note that although one can go from Case I to Case II and vice versa by the transformation $\sin\theta \longleftrightarrow \cos\theta\Theta_1$ and $\gamma_S \longleftrightarrow \gamma_1$, these cases are physically different. This is so because once θ and Θ_1 which parametrize the goldstino direction in the dilaton and the moduli VEV space are frozen, these cases will lead to different particle masses and will have physically distinct experimental consequences. Of course it is possible to view the two cases as part of a single case with a larger parameter space but we prefer to treat them as distinct on physical grounds. Again as in Case I we treat the phase of μ to be a free parameter and use the radiative breaking of the electro-weak symmetry to determine the magnitude of μ . An exhibition of the cancellation in EDMs for this case in the presence of large CP violating phases is given in Fig.7.

An interesting aspect of string models is that under the single scaling

$$m_{\frac{3}{2}} \rightarrow \lambda m_{\frac{3}{2}} \quad (53)$$

one has $F^S \rightarrow \lambda F^S$ and $F^i \rightarrow \lambda F^i$ and thus all the soft SUSY breaking parameters will have that scaling. We examine now the scaling phenomenon for the two models considered above. For this purpose it is useful to define $\lambda = \frac{m_{3/2}}{m_{3/2}^0}$ where $m_{3/2}$ is the running value and $m_{3/2}^0$ is $m_{3/2}$ at the extreme left. In Fig.8 we exhibit the result of the extrapolations for $\log_{10}\lambda^2|d_{e,n}|$ as a function of $m_{\frac{3}{2}}$ starting from a single point of cancellation at the far left. One finds that as $m_{\frac{3}{2}}$ increases the scaling is obeyed here to a good approximation. For comparison we also consider a heterotic string model. The cancellation for the EDMs for the type O-II model of Ref.[29] was discussed in Ref.[11]. We discuss here the scaling property. The soft SUSY breaking sector of this theory is parameterized by[29]

$$m_0^2 = \epsilon'(-\delta_{GS})m_{\frac{3}{2}}^2 \quad (54)$$

$$\tilde{m}_i = \sqrt{3}m_{\frac{3}{2}}(\sin\theta e^{-i\alpha_S} - \gamma_i\epsilon \cos\theta e^{-i\alpha_T}) \quad (55)$$

where $\gamma_1 = -\frac{33}{5} + \delta_{GS}$, $\gamma_2 = -1 + \delta_{GS}$, $\gamma_3 = 3 + \delta_{GS}$ and

$$A_0 = -\sqrt{3}m_{\frac{3}{2}}\sin\theta e^{-i\alpha_S} \quad (56)$$

The parameter δ_{GS} is fixed by the constraint of anomaly cancellation in a given orbifold model. The parameter μ and its phase are again treated as independent parameters. In Fig.9 we exhibit the result of the extrapolations for $\log_{10}\lambda^2|d_{e,n}|$ as a function of $\log_{10}m_{\frac{3}{2}}$ starting from a single point of cancellation at the far left. We find that scaling is obeyed for two of the three cases exhibited in Fig.9 over the entire range of $m_{\frac{3}{2}}$ considered. For the third case the initial part of the curves is in the non-scaling region and a new cancellation appears which, however, further reduces the EDM for this case maintaining consistency with the experimental EDM constraints. Eventually of course scaling seems to set in for this case as $m_{3/2}$ becomes larger. This third example is an interesting illustration of the approximate nature of the scaling analysis and of subleading non-scaling corrections. Since the cancellation is a rather delicate phenomenon these subleading terms can trigger a further cancellation which would lead to a breakdown of scaling. However, the EDM constraints are satisfied even more so in this case.

6 Conclusion

In this paper we discussed an algorithm for generating cancellations for the EDMs of the leptons and of the quarks within the framework of MSSM. We showed that

in theories where the μ parameter obeys the simple scaling behavior of Eq.(5) under the scaling of Eq.(3), the lepton and the quark EDMs show a simple scaling property in the $m_0 - m_{\frac{1}{2}}$ plane in the large μ region. Thus in this region the cancellation constraint on the electron and on the quark EDMs is essentially maintained under scaling. Thus given a single point in the SUSY parameter space in the large μ region where cancellations occur one can generate a trajectory in the $m_0 - m_{\frac{1}{2}}$ plane where cancellations are maintained by the use of scaling with only minor adjustments in other parameters. We emphasize that for low values of m_0 and $m_{\frac{1}{2}}$ some adjustment of the parameters to satisfy the EDM constraints will in general be needed to compensate for the fact that one is in the non-scaling region. We also discussed a class of Type IIB string models with 9 branes and 5₁ branes which have non-universal phases for the gaugino masses. We showed that such models can have large CP violating phases consistent with cancellations to guarantee the satisfaction of the EDM constraints. We also exhibited the existence of scaling in these models as well as in a heterotic string model. The simple algorithm described above with the caveats already discussed opens another window for the exploration of the SUSY parameter space with large CP phases and a relatively light SUSY particle spectrum. Finally as already pointed out in the second paper of Ref.[6] the cancellation hypothesis is an experimentally testable idea, i.e., that with soft SUSY phases $O(1 - 10^{-1})$ and with the SUSY spectrum within the naturalness limits of $O(1)$ TeV, the EDM of the electron and of the neutron should become visible with an order of magnitude improvement in the experimental EDM measurements. We further point out here that this observation is generic and should cover a range of models whether SUSY, string or brane. Such an order of magnitude improvement in experiment should be possible in the near future.

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7 Figure Captions

Fig.1. The trajectories in the $m_0 - m_{\frac{1}{2}}$ plane generated by scaling where cancellations occur in the SUSY EDMs consistent with the EDM constraints. (1) $|A_0| = 6.5$, $\theta_\mu = 2.92$, $\alpha_{A0} = -.4$, $\tan \beta = 4$, $\xi_1 = 0$, $\xi_2 = .2$ and $\xi_3 = .06$. (2) $|A_0| = 2.9$, $\theta_\mu = 3.02$, $\alpha_{A0} = .5$, $\tan \beta = 2.6$, $\xi_1 = .19$, $\xi_2 = .19$ and $\xi_3 = .41$. (3) $|A_0| = 5.5$, $\theta_\mu = 3.006$, $\alpha_{A0} = -.1$, $\tan \beta = 3.5$, $\xi_1 = .105$, $\xi_2 = .105$ and $\xi_3 = .15$. (4) $|A_0| = 4.4$, $\theta_\mu = 3.02$, $\alpha_{A0} = -.6$, $\tan \beta = 7$, $\xi_1 = 0$, $\xi_2 = .1$ and $\xi_3 = -.065$. (5) $|A_0| = 3.2$, $\theta_\mu = 2.8$, $\alpha_{A0} = -.4$, $\tan \beta = 5$, $\xi_1 = .31$, $\xi_2 = .3$ and $\xi_3 = .32$.

Fig.2. Plot of $\log_{10}|d_e|$ of the electron edm vs $m_{\frac{1}{2}}$ for the five cases of Fig.1.

Fig.3. Plot of $\log_{10}|d_n|$ of the neutron edm vs $m_{\frac{1}{2}}$ for the five cases of Fig.1.

Fig.4. Plot of $\log_{10}|\lambda^2 d_e|$ vs $m_{\frac{1}{2}}$ for the points in Fig. 2.

Fig.5. Plot of $\log_{10}|\lambda^2 d_n|$ vs $m_{\frac{1}{2}}$ for the points in Fig.3.

Fig.6. Plot of $\log_{10}|d_{e,n}|$ vs θ_μ for Model I based on 9 branes and 5_1 branes exhibiting the cancellation of the EDMs for the electron and for the neutron for the case when $m_{3/2} = 250$ GeV, $\theta = 1$, $\tan \beta = 5$, $\gamma_S = 0.295$, $\gamma_1 = 0.409$, $\theta_1 = 0.64$. The solid line is for the electron case and the dashed one is for the neutron.

Fig.7. Plot of $\log_{10}|d_{e,n}|$ vs θ_μ for Model II based on 9 branes and 5_1 branes exhibiting the cancellation of the EDMs for the electron and for the neutron for the case when $m_{3/2} = 500$ GeV, $\theta = 0.3$, $\tan \beta = 5$, $\gamma_S = 0.3$, $\gamma_1 = 0.4$, $\theta_1 = 0.9$. The solid line is for the electron case and the dashed one is for the neutron.

Fig.8. Plot of $\log_{10}\lambda^2|d_{e,n}|$ vs $m_{3/2}$ using one cancellation point (at far left) for each of the cases in Figs. 6 and 7 and scaling in $m_{3/2}$. θ_μ for each curve is fixed at the initial point to satisfy the experimental limits of edms by cancellation.

Fig.9. Plot of $\log_{10}\lambda^2|d_{e,n}|$ vs $\log_{10}m_{3/2}$ for the heterotic string model discussed in the text using one cancellation point (at far left) for each of the three cases. The parameters for the cases considered are: (1). $m_{3/2} = 1050$ GeV, $\theta = 0.06$, $\theta_\mu = 0.3$, $\tan \beta = 6$, $\alpha_S = 0.15$, $\alpha_T = 0.4$, $\delta_{GS} = -10$, $\epsilon = 0.006$, $\epsilon' = 0.001$, (2) $m_{3/2} = 340$ GeV, $\theta = 0.6$, $\theta_\mu = 0.3$, $\tan \beta = 3$, $\alpha_S = 0.25$, $\alpha_T = 0.37$, $\delta_{GS} = -4$, $\epsilon = 0.001$, $\epsilon' = 0.05$, (3) $m_{3/2} = 3$ TeV, $\theta = 0.05$, $\theta_\mu = 0.5$, $\tan \beta = 8$, $\alpha_S = 0.39$, $\alpha_T = 0.59$, $\delta_{GS} = -8$, $\epsilon = 0.004$, $\epsilon' = 0.0012$.

















